Loading a continuous-wave atom laser by optical pumping techniques

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(June 23, 2000)

Demonstrating that despite loss processes, Bose-Einstein condensates can be formed in steady state is a prerequisite for obtaining a coherent beam of atoms in a continuous-wave atom laser. In this paper we propose a method for loading atoms into the thermal component of a Bose condensed cloud confined in a magnetic trap. This method is aimed at allowing steady state dynamics to be achieved. The proposed scheme involves loading atoms into the conservative magnetic potential using the spontaneous emission of photons. We show that the probability for the reabsorption of these photons may be small .

03.75.Fi, 05.20.Dd, 32.80.Pj

I. INTRODUCTION

An atom laser is a device which generates an atomic beam which is both intense and coherent [1–10]. It is typically defined by direct analogy with the optical laser. As in the case of an optical laser, if a threshold condition is satisfied the atom laser will operate far from equilibrium in a dynamical steady state [11]. In this regime, the macroscopically occupied quantum state is continuously depleted by loss to the output field, and continuously replenished by pumping from an active medium. The threshold condition for the atom laser is analogous to the critical point associated with the phase transition from a normal gas to a Bose-Einstein condensate (BEC). A version of the "pulsed" atom laser [12,13] was demonstrated soon after the first experimental realizations of BEC in dilute alkali gases [14–17].

The three basic components of an atom laser are: a cavity (or resonator), an active medium (or reservoir), and an output coupler. The cavity is typically a three-dimensional magnetic trap which for containment uses the force on the atomic magnetic dipole via an inhomogeneous potential. The active medium may be composed of a saturated dilute thermal gas formed through evaporative cooling or other techniques. Saturated in this context implies that the energy distribution function of the atoms is such that thermodynamic relaxation of the gas through binary collisions would result in the stimulated accumulation of atoms into the ground state. The output coupler can be implemented by either applying a short radio-frequency (RF) pulse to the condensate and thereby separating the condensate into a trapped component and an un-trapped component [12,13,18,19], or alternatively by creating a situation in which the atoms tunnel out through a potential barrier [20]. A recent demonstration of a quasi-continuous atom laser used an optical Raman pulse to drive transitions between trapped and un-trapped magnetic sublevels giving the output-coupled BEC fraction a non-zero momentum [21].

To date, no continuous wave (CW) atom laser has been demonstrated in which there is replenishing of the reservoir in direct analogy with CW-optical laser. Replenishing the reservoir is required in steady state to compensate for the loss due to output coupling as well as to compensate for various intrinsic loss mechanisms such as collisions with hot atoms from the background vapor and inelastic two-body and three-body collisions between trapped atoms [1,10,22].

In this paper we propose a scheme for loading the thermal cloud into a magnetic trap by optically pumping atoms from an external cold atomic beam source. Optical pumping is necessary since the magnetic potential is conservative and therefore a dissipative process is required for confinement of the atoms to be achieved. When combined with evaporative cooling, our approach may allow both steady state BEC and the CW atom laser to be demonstrated. The crucial point is that while one is trying to obtain a steady state BEC by injecting atoms into the magnetic trap, the loading mechanism itself should not heat the system or destroy the delicate condensate component.

There have been related proposals to ours using single atom "dark states" for reaching low temperatures by purely optical means—even well below the photon recoil limit [23–25]. As shown by Castin *et al.* in Ref. [26], the main hurdle in achieving BEC in all these schemes is that of photon reabsorption. Each reabsorption can remove one atom from the dark state and increment the energy of the cloud on average by one recoil energy. We engineer the analogous effect in our scheme to be negligible by configuring the frequency of the emitted photons and adjusting the spatial geometry of the system to reduce the probability of reabsorption by the condensate.

II. THE MODEL

We consider atoms whose internal degrees of freedom are described by a three level atomic system as shown in Fig. 1. Our description of the loading scheme is motivated by the experimental setup in Ref. [27]. There the atoms were loaded from a background vapor and precooled in a magneto-optical trap (MOT). They were then launched into the second MOT, optically pumped to the

desired internal atomic hyperfine state, and transfered to a magnetically confining potential in which evaporative cooling was implemented. Here we study the possibility for loading the thermal cloud into the magnetic trap directly and continuously. This would allow the evaporative cooling to be carried out in steady state, which is the situation considered in Ref. [11].

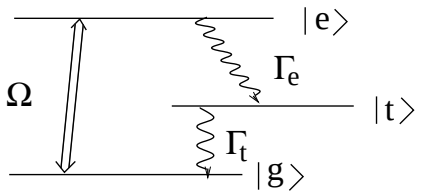

FIG. 1. Level scheme illustrating: the ground state $|g\rangle$, excited state $|e\rangle$, and trap state $|t\rangle$. A laser is used to couple states $|g\rangle$ and $|e\rangle$ with an intensity characterized by the Rabi frequency Ω. Spontaneous emission of photons occurs with rate Γ_e from $|e\rangle$ to $|t\rangle$ and rate Γ_t from $|t\rangle$ to $|g\rangle$.

Our scheme is illustrated in Fig. 2. Atoms are injected into the trap in state $|g\rangle$. As the atoms enter the spatial region containing the thermal cloud they pass through a laser field and are coherently pumped to the excited state $|e\rangle$ from which they may spontaneously emit a photon and end up in the trap state $|t\rangle$.

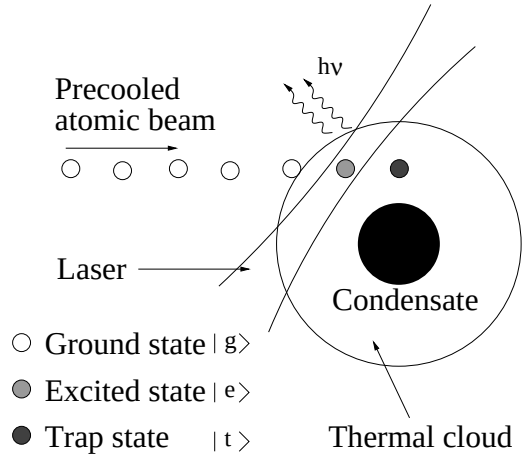

FIG. 2. Schematic diagram showing the proposed loading scheme for obtaining a steady state BEC and CW atom laser. Atoms entering the laser field are in the ground state. As the atoms pass through the laser field, the population is continuously transferred to the trap state by a sequence of coherent pumping and spontaneous emission.

As will be shown later the spontaneously emitted photons may be out of resonance with the bare $|t\rangle \rightarrow |e\rangle$ transition and therefore reabsorption by atoms in the thermal cloud is greatly reduced (the exception to this general statement is the negligible fraction of thermal atoms in the localized spatial region of the focussed laser beam.) Using this approach atoms may be continuously loaded into the thermal cloud and be confined by the trapping potential without prohibitive reabsorption of the spontaneous photons emitted in the required dissipative process.

We use a master equation [28,29] to study the evolution of this system. We define the spontaneous emission rates between the excited state and trap state and between the trap state and ground state as Γ_e and Γ_t respectively, and for simplicity neglect spontaneous emission directly from $|e\rangle$ to $|g\rangle$. Assuming that the laser is on resonance with the atomic transition frequency (detuning $\delta = 0$), the coherent interaction can be modeled by the Hamiltonian in the interaction picture

$$H = \frac{\Omega}{2}\hbar(\sigma_+^{ge} + \sigma_-^{ge}) \tag{1}$$

where $\sigma_+^{ge}, \sigma_-^{ge}$ are the raising and lowering operators respectively between the ground state $|g\rangle$ and the excited state $|e\rangle$. The master equation for the reduced density operator of the atom includes terms due to both the spontaneous emission and a part corresponding to the coherent driving source. In the interaction picture this takes the form

$$\begin{aligned}\frac{\partial\rho}{\partial t} = &-i\frac{\Omega}{2}[\sigma_+^{ge} + \sigma_-^{ge}, \rho] \\ &+ \frac{\Gamma_e}{2}(2\sigma_-^{et}\rho\sigma_+^{et} - \sigma_+^{et}\sigma_-^{et}\rho - \rho\sigma_+^{et}\sigma_-^{et}) \\ &+ \frac{\Gamma_t}{2}(2\sigma_-^{gt}\rho\sigma_+^{gt} - \sigma_+^{gt}\sigma_-^{gt}\rho - \rho\sigma_+^{gt}\sigma_-^{gt}),\end{aligned} \tag{2}$$

where $\sigma_+^{et}, \sigma_-^{et}$ are the raising and lowering operators respectively between the excited and trap state and $\sigma_+^{gt}, \sigma_-^{gt}$ are the raising and lowering operators respectively between the trap state and ground state. From the above master equation we obtain the coupled differential equations for the evolution of matrix elements of ρ with the constraint $\mathrm{Tr}[\rho] = 1$. The atomic population in the three levels given by the matrix elements $\rho_{gg}(t), \rho_{ee}(t)$ and $\rho_{tt}(t)$ are obtained for values of $\Gamma_e = 1$ and $\Omega = 5$ and initial condition given by $\rho_{gg}(0) = 1$, $\rho_{ee}(0) = \rho_{tt}(0) = 0$. We consider two different values for Γ_t; $\Gamma_t = 2$ and $\Gamma_t = 0$, which give qualitatively distinct dynamics. The zero value for Γ_t corresponds to optically pumping the atom from the ground to trap state, and leads to complete transfer of population at long times. The non-zero value of Γ_t gives a steady-state solution for the level populations in which the atom cycles indefinitely, which will be useful later for the calculation of the spectrum of the spontaneously emitted photons.

The steady state spectrum of spontaneously emitted photons on the $|e\rangle \rightarrow |t\rangle$ transition is given by the Fourier transform of the two-time correlation function

$$K(t+\tau, t) = \langle\sigma_+^{et}(t+\tau)\sigma_-^{et}(t)\rangle, \tag{3}$$

the differential equation for which is obtained by combining the equation for ρ_{tg} and ρ_{te} above along with the quantum regression theorem (QRT) [29].

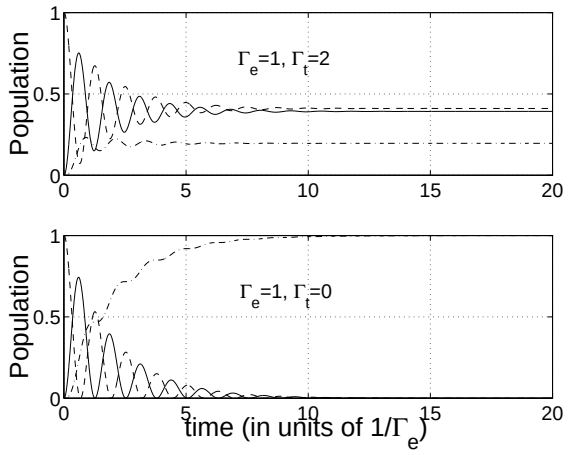

FIG. 3. Populations as a function of time; ground state $|g\rangle$ shown by a solid line, excited state $|e\rangle$ shown by dashed and trap state $|t\rangle$ shown by dot-dashed line. The parameters used were $\Omega = 5, \Gamma_e = 1$ and (a) $\Gamma_t = 2$, (b) $\Gamma_t = 0$.

The QRT states that for a set of operators Y_i, which evolve in steady-state according to

$$\frac{\partial}{\partial\tau}\langle Y_i(\tau)\rangle = \sum_j G_{ij}\langle Y_j(\tau)\rangle, \tag{4}$$

then two-time averages can be related to the one-time expectation values by

$$\frac{\partial}{\partial\tau}\langle Y_i(t+\tau)Y_l(t)\rangle = \sum_j G_{ij}\langle Y_j(t+\tau)Y_l(t)\rangle, \tag{5}$$

where the G_{ij}'s are the same coefficients in both Eq. (4) and Eq. (5). For a general non-steady-state problem in which transient dynamics must be considered, application of the QRT is typically more complicated. For simplicity our approach is to take a fictitious non-zero value of Γ_t, and then consider the limit $\Gamma_t \to 0$. The spectrum defined as

$$S(\omega) = \lim_{t\to\infty}\int_{-t}^{t} d\tau K(\tau)e^{-i\omega\tau} \tag{6}$$

is then given by

$$S(\omega) = \frac{\rho_{tt}^{ss}}{\pi\Omega_{\rm eff}}\left[\frac{\left(i\frac{\Gamma_e}{4}+\frac{\Omega_{\rm eff}}{2}\right)\left(\frac{\Gamma_e}{4}+\frac{\Gamma_t}{2}\right)}{\left(\frac{\Gamma_e}{4}+\frac{\Gamma_t}{2}\right)^2+\left(\omega+\frac{\Omega_{\rm eff}}{2}\right)^2}\right] + \frac{\rho_{tt}^{ss}}{\pi\Omega_{\rm eff}}\left[\frac{\left(-i\frac{\Gamma_e}{4}+\frac{\Omega_{\rm eff}}{2}\right)\left(\frac{\Gamma_e}{4}+\frac{\Gamma_t}{2}\right)}{\left(\frac{\Gamma_e}{4}+\frac{\Gamma_t}{2}\right)^2+\left(\omega-\frac{\Omega_{\rm eff}}{2}\right)^2}\right] \tag{7}$$

where $\Omega_{\rm eff} = \sqrt{\Omega^2 - \Gamma_e^2/4}$

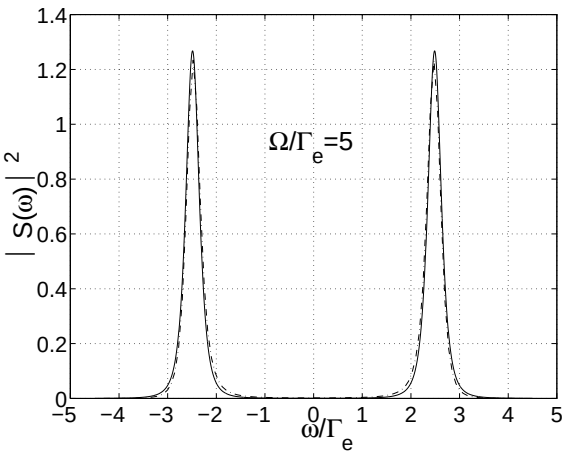

FIG. 4. Comparison of the $|e\rangle \to |t\rangle$ transition spectrum using the steady state calculation (solid) and the quantum trajectory method (dash-dot).

In Fig. 4 we plot $|S(\omega)|^2$ shown by the solid curve for the limiting case of $\Gamma_t \to 0$. The spectrum is AC-Stark split into two Lorentzians separated by $\Omega_{\rm eff}$. This spectrum is commonly referred to as the Autler-Townes doublet [31,32].

For comparison we now use the method described in Ref. [30] to provide an alternative calculation of the fluorescence spectrum, extending our treatment to include the possibility for the time dependence of Ω. This involves the decomposition into "quantum trajectories", where the evolution is conditional on a certain record of dissipative events. In the case of the evolution conditional on the spontaneous emission of no photons, the quantum trajectory $|\psi(t)\rangle$, evolves according to

$$\frac{\partial|\psi(t)\rangle}{\partial t} = \frac{1}{i\hbar}H_{\rm eff}|\psi(t)\rangle, \tag{8}$$

where $H_{\rm eff}$ is the non-Hermitian Hamiltonian defined by

$$H_{\rm eff} = H - i\frac{\hbar\Gamma_e}{2}\sigma_+^{et}\sigma_-^{et}. \tag{9}$$

For the evolution conditional on the spontaneous emission of one photon, each quantum trajectory $|\psi_{\omega_j}(t)\rangle$ is labelled by the frequency of the photon ω_j. If we consider a time interval $t \in [0,\tau]$, then according to the Fourier sampling theorem a complete description requires choosing each ω_j from a discrete but infinite set of frequencies spaced $(2\pi/\tau)$ apart. The observed spectrum may then be defined as

$$S(\omega) = \langle\psi_\omega(\tau)|\psi_\omega(\tau)\rangle \tag{10}$$

with (see Ref. [30])

$$\frac{d}{dt}|\psi_{\omega_j}(t)\rangle = \sqrt{\frac{\Gamma_e}{\tau}}\sigma_-^{et}|\psi(t)\rangle + \frac{1}{i\hbar}(H_{\rm eff}+\hbar\omega_j)|\psi_{\omega_j}(t)\rangle. \tag{11}$$

By solving the coupled Eqns. (8) and (11) we get

$$|\psi_{\omega_j}(t)\rangle \;=\; \begin{pmatrix} 0 \\ 0 \\ f_{\omega_j}(t) \end{pmatrix} \tag{12}$$

where

$$\begin{aligned} f_{\omega_j}(t) &= q\mathcal{M}N^{-1}\left((i\omega_j-\Gamma_e/4)^2+\Omega_{\rm eff}^2/4\right)^{-1} \\ \mathcal{M} &= \sin\left(\frac{\Omega_{\rm eff}t}{2}\right)e^{-\Gamma_e t/4}(i\omega_j-\Gamma_e/4) \\ &\quad -\frac{\Omega_{\rm eff}}{2}\cos\left(\frac{\Omega_{\rm eff}t}{2}\right)e^{-\Gamma_e t/4}+\frac{\Omega_{\rm eff}}{2}e^{-i\omega_j t} \\ q &= \frac{4}{\Omega}\sqrt{\frac{\Gamma_e}{\tau}}\left[\frac{\lambda_+\lambda_-}{\lambda_+-\lambda_-}\right] \end{aligned} \tag{13}$$

The single-photon spectrum which is so obtained, $|S(\omega)|^2$, is shown by dot-dashed curve in Fig. 4. The normalization constant N is introduced so that $|f_{\omega_j}|^2 \to 1$ as $t \to \infty$. We see a close agreement between the spectra calculated by the two methods, the quantum regression theorem and the quantum trajectory approach of solving conditional dynamics. The crucial observation is the absence of a significant contribution in the zero-frequency region of the fluorescence spectrum where there would be strong photon reabsorption by the thermal cloud. Virtually all of the spontaneous photons are off resonance with the bare $|t\rangle \to |e\rangle$ transition which is the resonant frequency for all thermal atoms out of the laser field focus.

III. EFFECT OF TIME DEPENDENT Ω AND FINITE DETUNING ON THE TRANSIENT SPECTRUM

The assumed time independent nature of the interaction picture Hamiltonian (1) may not be a realistic situation. In an actual experiment, as the atoms pass through the laser beam they would see a continuously changing intensity and hence a changing Ω. At the same time, existence of finite detuning δ (assumed previously to be zero for simplicity), may result in the time dependence of the operators σ_+^{ge} and σ_-^{ge} given by

$$\sigma_\pm^{ge}(t) = \sigma_\pm^{ge} e^{\pm i\delta t} \tag{14}$$

Thus, in a more realistic situation the interaction picture Hamiltonian of Eq.(1) should be replaced by the time dependent form

$$H(t) = \frac{\Omega(t)}{2}\hbar(\sigma_+^{ge}e^{i\delta t}+\sigma_-^{ge}e^{-i\delta t}) \tag{15}$$

Using this interaction Hamiltonian we now recalculate the AC-Stark split spectrum. Firstly we consider the $\delta = 0$ case with the time dependence solely due to $\Omega(t)$. Since the problem is a transient one in which population is completely transferred to the trap state over a certain time scale, the variation of Ω over this time scale is crucial. Also, from the calculations in the previous section it is clear that a large value of Ω is favorable for our scheme. We therefore consider various time scales over which Ω reaches its maximum value compared to the optical pumping time as shown in Fig. 5. The corresponding spectrum shown in Fig. 6 is calculated using the method of conditional dynamics discussed in the previous section. The calculation shows that the crucial requirement for our model i.e. that the emitted photons cannot be resonantly absorbed by the thermal cloud outside the laser focus,is only satisfied if the atom sees a field increasing in intensity on a time scale which is short compared to the time it takes to optically pump the population from $|g\rangle$ to $|t\rangle$. For a particular average velocity of the atoms in the atom beam source, this gives the requirement for the spatial variation of the laser intensity associated with the beam waist. Alternatively, the laser field could be pulsed on and off to achieve a sufficiently rapid rise time.

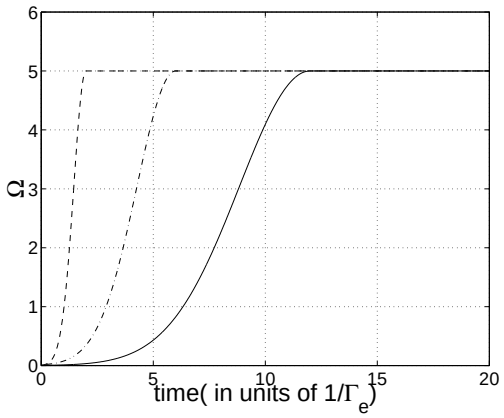

FIG. 5. Ω as a function of time, t. The three different curves correspond to different time scales over which Ω reaches its maximum value.

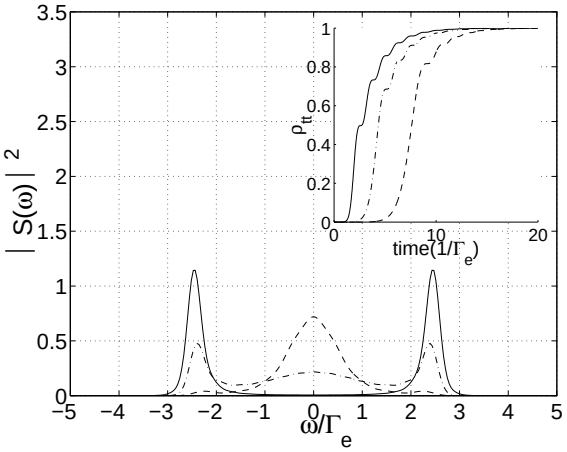

FIG. 6. Transient spectrum for $|e\rangle \to |t\rangle$ transition for the different scenarios for $\Omega(t)$ shown in Fig.5.

Now we consider the effect of finite detuning, δ, in addition to the time dependence of Ω. For the case of constant Ω, a finite δ has the effect of transforming Ω to $\sqrt{\Omega^2+\delta^2}$ [32]. Therefore, even if Ω depends on time, it is intuitive to expect a similar effect resulting in an increased AC-Stark splitting. This is illustrated in Fig. 7 where the spectrum is calculated for the case of $\Omega(t)$ shown by dot-dashed curve of Fig. 5 for various values of δ. Thus a large detuning may seem to be favorable for our purpose since the spontaneous photons are emitted further away from resonance. On the other hand, large detuning means slower optical pumping rate and hence a wider laser focus will be required for the same $\Omega(t)$ as illustrated in the inset of Fig. 7. Therefore, it is important to achieve an optimum set of parameters that would allow a maximum possible splitting depending on the average atomic speed and line width of the atomic levels under consideration.

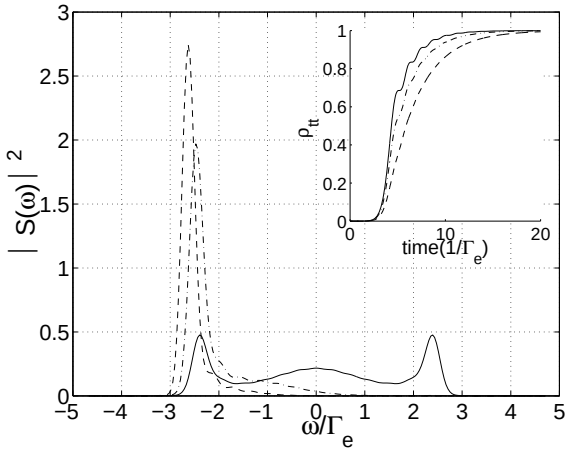

FIG. 7. Figure shows the effect of detuning on the transient spectrum for $|e\rangle \rightarrow |t\rangle$ transition. Dashed for $\delta = 0$, solid for $\delta = 1$ and dot-dashed for $\delta = 2$.

IV. CONCLUSION

The AC-Stark splitting of the spectrum indicates that this scheme may potentially be used to provide a method for continuous loading of the thermal cloud in an atom laser. In the case of constant Ω the spectrum of the spontaneously emitted photons consists of two peaks separated by $\Omega_{\rm eff}$. Even for the realistic case of time varying Ω, large splitting between the two peaks can be achieved provided that the rise time of the laser field intensity is sufficiently small. The splitting can be further enhanced by introducing a finite detuning.

Large splitting of the spectrum is essential to avoid heating due to inelastic multiple scattering of photons. Usually heating in an optical pumping scheme is caused by photon reabsorption and re-emission. Each reabsorption can remove one atom from the tr ap state and increment the energy of the cloud on average by one recoil energy. This effect is significantly reduced in our model since the spontaneously emitted photons are off-resonance with the $|t\rangle \rightarrow |e\rangle$ transition and cannot be reabsorbed by atoms outside the laser focus. Therefore the heating of the thermal component due to photon reabsorption may be suppressed. One can also reduce the cross section for photon reabsorption by designing the shape of the trap, for example a cigar shape or disc shape, such that there is a limited solid angle for photon reabsorption. Another possibility is to align the optical pumping laser and the geometry of the trap to use the intrinsic dipole radiation pattern to reduce spontaneous emission into unfavorable directions.

Our model requires that the atoms be in the trap state before they are out of the laser field, which may place a limit on the maximum speed of atoms in the atomic beam source. However, the faster the incoming atoms, the smaller the rise time of Ω which as previously discussed is required to avoid the emission of photons at zero-frequency. The optimum speed of incoming atoms is therefore a balance between these two requirements. Since our goal is to obtain BEC in a steady state and not just BEC, one possibility to obtain sufficiently slow speeds is to use a condensate prepared in another trap as a source instead of just laser precooled atoms formed from a magneto-optical-trap.

For a typical alkali atom, the Zeeman level diagram is composed of many hyperfine levels. Even though our three level system is a simplification of a real situation, most of the omitted processes, such as imperfect optical pumping to the trapping state, will lead to additional loss of atoms and modify the effective incident flux. This will impact the possibility for achieving a sufficiently high pumping rate of cold atoms to allow the CW atom laser and steady state BEC to be achieved, as discussed in Ref. [11].

We would like to thank Eric Cornell, John Cooper and R. Walser for discussions. We acknowledge support from the U.S. Department of Energy, Office of Basic Energy Sciences via the Chemical Sciences, Geosciences and Biosciences Division.

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